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# Heavy Quark Contributions to $W$ Pair Production at Hadron Colliders

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## Abstract

We investigate the effects of including fermion masses in calculating the $W$ pair polarization density matrix for the process $q\bar{q} \to W^+W^-$, and in testing the neutral scalar and vector boson couplings to $W^+W^-$ at hadron colliders. Appreciable effects are found as SSC energy and, in some cases, at Tevatron energy as well.

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$W$ pair production has attracted a lot of attention [1-20] because it provides a channel for testing the gauge couplings of the electroweak vector bosons and for probing the neutral scalar and vector boson sectors of a given model. As far as testing the gauge couplings of the standard electroweak theory is concerned, the electron-positron colliders are perhaps the better machines. Probing the neutral scalar or vector boson sectors, however, is not the case. For one thing, the hadron colliders have higher reach for the same cost. For the other, the neutral scalar couplings to the fermions are usually proportional to their masses. The smallness of the electron mass relative to the $W$ mass makes the direct probing of these neutral scalars through $W$ pair production at $e^+e^-$ colliders impossible. Similarly, the general neutral vector boson couplings to $W^+W^-$ may contain form factors which are impossible to probe at $e^+e^-$ colliders. To take advantage of these effects at hadron colliders, it is necessary to investigate the contributions to $W$ pair production from heavy quark annihilation. However, in all the previous discussions on $W$ production at hadron colliders [3,6,7-10,14,17,20], the fermion masses were ignored. As we shall see, this is justified even up to SSC energy if one is interested *only* in the total cross section or the polarization-summed differential cross section of $W$ pair productions away from s-channel resonances and in the context of the Standard Model.

In this note, we will investigate the effects of including the quark masses in calculating the polarization density matrix of the $W$ pairs. We are particularly interested in the roles of neutral scalar and vector bosons with general couplings to the $W$'s consistent with Lorentz invariance. When quark masses are taken into consideration, there are, in general, three kinds of effects. First, there are additional terms to the polarization density matrix corresponding to the cases when the incoming quark and antiquark have the same helicity. For non-standard model couplings of neutral scalar or vector bosons, these terms may violate tree level unitarity and increase with increasing center-of-mass energy despite the fact that they are proportional to the mass of the incoming fermion. The luminosities of heavy quarks, however, are small and whether the additional terms are numerically important has to be examined carefully. Second, the cancellations between $s$- and $t$-channel exchange diagrams characteristic of gauge theories are less effective when the $t$-channel quark masses are taken into account. Finally, the kinematics of the heavy quarks may affect their effective luminosities. Consider the case when the top quark mass is close to the $W$ mass. Then one or both of the incoming top-quarks may carry a tiny fraction of the nucleon momentum and the C.M. energy is still above the $W$ pair production threshold. Since the luminosity of sea quarks with smaller $x$—the fraction of nucleon momentum they carry—are larger, taking into account the heavy quark mass has the effect of decreasing $x$ and hence increasing their effective luminosities. This also pushes us into the small $x$ region of quark distribution functions. In our numerical works, we use the parton distribution function of set 2 of EHLQ [20]. We take the top quark mass $m_t = 70 GeV$ and to avoid the region $x < 10^{-4}$ which is beyond the limit of validity of the EHLQ structure functions, we make the conservative cut by requiring that $x$ be greater than $\tau = \hat{s}/s$ when integrating over the rapidity of the $W$ pair, ($\hat{s}$ is the parton C.M. energy squared and $s$ is the corresponding quantity for incoming nucleons) which is satisfied automatically if quark masses are neglected. One should keep in mind, however, that there is still no complete understanding as to how to deal with the heavy quark distribution functions in a consistent way [20,21].

Let $M(k_1\sigma_1, k_2\sigma_2; q_1\lambda_1, q_2\lambda_2)$ be the scattering amplitude for a quark and an anti-quark with momenta $k_1, k_2$ and helicities $\sigma_1, \sigma_2$ respectively to annihilate into $W^-W^+$ with momenta $q_1, q_2$ and polarizations $\lambda_1, \lambda_2$ respectively. We use rectangular basis for the polarizations of $W$ and follow the notation of Ref. [19], so that $\lambda = 1$ labels the transverse polarization in the production plane, $\lambda = 2$ labels the transverse polarization perpendicular to the production plane, and $\lambda = 3$ labels the longitudinal polarization. The polarization density matrix is defined by

$$\mathcal{P}_{\lambda_1\lambda_2;\lambda_1'\lambda_2'} = \sum_{\sigma_1,\sigma_2} M\left(k_1\sigma_1, k_2\sigma_2; q_1\lambda_1, q_2\lambda_2\right) M^*\left(k_1\sigma_1, k_2\sigma_2; q_1\lambda_1', q_2\lambda_2'\right) \qquad (.1)$$

and can be written as

$$\mathcal{P} = 2e^4\tilde{\mathcal{P}}, \quad \tilde{\mathcal{P}} = R^v \otimes R^v + R^a \otimes R^a + S^v \otimes S^v + S^a \otimes S^a \qquad (.2)$$

The four terms are related to four initial polarization states of the quark and anti-quark. For massless quarks, only $R^v$ and $R^a$ are non-zero. They have been given in Ref. [4,19] for general three vector boson couplings (with on shell $W^-W^+$)

$$\Gamma_V^{\alpha\beta\mu}(q_1, q_2, p) = f_1^V (q_1 - q_2)^\mu g^{\alpha\beta} - \frac{f_2^V}{m_w^2}(q_1 - q_2)^\mu p^\alpha p^\beta + f_3^V \left(p^\alpha g^{\mu\beta} - p^\beta g^{\mu\alpha}\right)$$

$$
\begin{aligned}
&+ \quad i f_4^V \left(p^\alpha g^{\mu\beta} + p^\beta g^{\mu\alpha}\right) + i f_5^V \varepsilon^{\mu\alpha\beta\rho} \left(q_1 - q_2\right)_\rho \\
&- \quad f_6^V \varepsilon^{\mu\alpha\beta\rho} p_\rho - \frac{f_7^V}{m_w^2} \left(q_1 - q_2\right)^\mu \varepsilon^{\alpha\beta\rho\sigma} p_\rho \left(q_1 - q_2\right)_\sigma \\
&+ \quad i f_8^V p^\mu g^{\alpha\beta} - i \frac{f_9^V}{m_w^2} p^\mu p^\alpha p^\beta - i \frac{f_{10}^V}{m_w^2} p^\mu \varepsilon^{\alpha\beta\rho\sigma} p_\rho \left(q_1 - q_2\right)_\sigma
\end{aligned}
\tag{.3}
$$

where $p = q_1 + q_2$.

The last three form factors give contributions proportional to the mass of incoming fermions and are neglected in Ref. [4,19]. We shall focus on the effects of these last three form factors in hadron colliders in this note while the effects of the other form factors will be discussed elsewhere [22]. In Standard Model, $f_1^\gamma = f_1^Z = 1$ and $f_3^\gamma = f_3^Z = 2$ are the only non-vanishing form factors at tree level. Additional neutral vector bosons with more general couplings are predicted by various extensions beyond Standard Model such as left–right symmetric [23], technicolor [24], composite [25] models, etc.

We also investigate the effects of a general neutral scalar–vector boson coupling

$$
\Gamma_S^{\alpha\beta} = f_1^S g^{\alpha\beta} - \frac{f_2^S}{m_w^2} p^\alpha p^\beta - \frac{f_3^S}{m_w^2} \varepsilon^{\alpha\beta\rho\sigma} p_\rho \left(q_1 - q_2\right)_\sigma . \tag{.4}
$$

For Standard Model Higgs, $f_1^S = 1, f_2^S = f_3^S = 0$. We allow neutral scalars with $\gamma_5$ coupling to the fermions as well. The couplings are assumed to be proportional to the fermion mass as the standard Higgs. For this type of neutral scalars, we consider again a general coupling to vector bosons

$$
\Gamma_P^{\alpha\beta} = i f_1^P g^{\alpha\beta} - i \frac{f_2^P}{m_w^2} p^\alpha p^\beta - i \frac{f_3^P}{m_w^2} \varepsilon^{\alpha\beta\rho\sigma} p_\rho (q_1 - q_2)_\sigma . \tag{.5}
$$

Neutral scalars with such couplings are predicted, for example, in extended technicolor models [26]. Only the following combinations of form factors appear in the polarization density matrix

$$
a_n = Q_i \left(f_n^\gamma - \frac{\hat{s}}{\hat{s} - m_Z^2} f_n^Z\right), \; b_n = -\left(\frac{L_i - R_i}{2\sin^2\theta_w} \frac{\hat{s}}{\hat{s} - m_Z^2} f_n^Z\right), n = 1, \cdots 10 \tag{.6}
$$

$$
b = \frac{1}{4\sin^2\theta_w} \sum_I \frac{\hat{s} \mid U_{iI} \mid^2}{\hat{t} - m_I^2}, \;\; b_n^{S,P} = \frac{1}{2\sin^2\theta_w} \; \frac{\hat{s}}{\hat{s} - m_H^2 + i m_H \Gamma_H} f_n^{S,P}, \;\; n = 1, 2, 3 . \tag{.7}
$$

In the above formulas, $Q_i$ are the electric charges of the incoming quarks, $L_i, R_i$ their left and right weak charges (coupling strength to $Z$); $m_I$ are the masses of the $t$-channel quarks, $U_{iI}$ are the quark mixing matrix elements; $m_H$ and $\Gamma_H$ are the masses and widths of the neutral scalars.

The matrices $R^v, R^a, S^v, S^a$ in equation (2) satisfy the following relations if we treat $a_n, b_n, b, b_n^{S,P}$ as real quantities [19]:

$$A_{12} = A_{21}^*, \quad A_{13} = -A_{31}^*, \quad A_{23} = -A_{32}^* \tag{.8}$$

where $A$ represents $R^v, R^a, S^v$ or $S^a$ and '$*$' means complex conjugate. The matrix elements of $S^v, S^a$ for down-type incoming quarks are given by

$$\begin{aligned}
S_{11}^v &= \frac{1}{2\gamma_i}\left[-\beta_i b_1^S - (2a_1 - b_1)\beta_w\cos\theta - b(\beta_i - \beta_w\cos\theta) + 2b\beta_i^2\sin^2\theta\right] \\
S_{12}^v &= \frac{1}{2\gamma_i}\left\{\left[2a_6 - b_6 + 4(2a_7 - b_7)\beta_w^2\gamma_w^2\right]\cos\theta + 4b_3^S\beta_i\beta_w\gamma_w^2\right\} \\
S_{13}^v &= -\frac{\gamma_w}{2\gamma_i}\left[(2a_- - b_-)\beta_w - 2b(\beta_w - \beta_i\cos\theta) + b\beta_w\right]\sin\theta \\
S_{22}^v &= \frac{1}{2\gamma_i}\left[\beta_i b_1^S + (2a_1 - b_1)\beta_w\cos\theta + b(\beta_i - \beta_w\cos\theta)\right] \\
S_{23}^v &= -\frac{i\gamma_w}{2\gamma_i}\left[(2a_5 - b_5)\beta_w^2 + i(2a_6 - b_6) - b\beta_w^2\right]\sin\theta \\
S_{33}^v &= \frac{1}{2\gamma_i}\left\{\beta_i\left[b_1^S(2\gamma_w^2 - 1) - 4b_2^S\beta_w^2\gamma_w^4\right] + \left[(2a_1 - b_1)(2\gamma_w^2 - 1) - 4(2a_2 - b_2)\beta_w^2\gamma_w^4\right]\beta_w\cos\theta\right. \\
&- \left. 2(2a_3 - b_3)\gamma_w^2\beta_w\cos\theta + b(2\gamma_w^2 - 1)(\beta_i - \beta_w\cos\theta) + 2b\gamma_w^2(\beta_w - \beta_i\cos\theta)\cos\theta\right\}
\end{aligned} \tag{.9}$$

$$\begin{aligned}
S_{11}^a &= -\frac{1}{2\gamma_i}\left[\left(4\gamma_z^2 - 1\right)b_8 + b_1^P\right] \\
S_{12}^a &= \frac{1}{2\gamma_i}\left[\left(4\beta_w\gamma_w^2 b_{10} - b_5\beta_w\right)\left(4\gamma_z^2 - 1\right) + 4\beta_w\gamma_w^2 b_3^P + b\left(\beta_w - \beta_i\cos\theta\right)\right] \\
S_{13}^a &= -\frac{i\gamma_w}{2\gamma_i}b\beta_i\beta_w\sin\theta \\
S_{22}^2 &= -S_{11}^a \\
S_{23}^a &= -\frac{\gamma_w}{2\gamma_i}b\beta_i\sin\theta
\end{aligned} \tag{.10}$$

$$S^a_{33} = \frac{1}{2\gamma_i}\left\{\left[b_8(2\gamma_w^2-1)-4b_9\beta_w^2\gamma_w^4+2b_4\beta_w^2\gamma_w^2\right](4\gamma_z^2-1)+b_1^P(2\gamma_w^2-1)-4b_2^P\beta_w^2\gamma_w^4\right\}$$

where $\gamma_i, \gamma_w, \gamma_z, \beta_i, \beta_w, \beta_z$ are, as usual,

$$\gamma_w = \frac{\sqrt{\hat{s}}}{2m_w} \quad \beta_w = \left(1-\frac{4m_w^2}{\hat{s}}\right)^{\frac{1}{2}} \text{ etc.} \tag{.11}$$

and $a_- = a_3 - ia_4, b_- = b_3 - ib_4$. In these formulas, $\theta$ is the angle between $W^-$ and the incoming quark in the parton $CM$ frame and the indices $i,j$ in $S^{v,a}_{i,j}$ refer to the polarizations of $W^-$ and $W^+$ respectively.

For up-type quarks, we make the change $\theta \to \pi - \theta$ in equations (9) and (10) (which also affect $b$ through the $t$-channel momentum square $t$), take the transpose of the matrices and substitute the up-type quark masses and charges in equations (6) and (7). The signs of $f_n^v$ for $n = 1,2,3,6,7$ also have to be flipped.

Let $M$ be the invariant mass of the $W$ pair. Then we have

$$\left[\frac{d\sigma}{dMd\cos\theta^*}\right]_{\lambda_1\lambda_2,\lambda_1'\lambda_2'} = \frac{\pi\alpha^2\beta_w}{s\beta_i M N_c}\sum_i\int dy_+\frac{1}{\beta_i}\left(\cos h^2 y_+ - \beta_i^2\sin h^2 y_+\right) \cdot$$

$$\left[f_i^A\left(x_a,M^2\right)f_{\bar{i}}^B\left(x_b,M^2\right)\tilde{P}\left(\theta^*,M\right)+f_{\bar{i}}^A\left(x_a,M^2\right)f_i^B\left(x_b,M^2\right)\tilde{P}\left(\theta^*+\pi,M\right)\right]_{\lambda_1\lambda_2,\lambda_1'\lambda_2'} \tag{.12}$$

where $y_+ = \frac{1}{2}(y_1+y_2), N_c = 3, \theta^*$ is related to the rapidities $y_1, y_2$ of $W^-, W^+$ by

$$\beta_w\cos\theta^* = \tanh\frac{1}{2}(y_1-y_2) \tag{.13}$$

and the summation runs through quark flavors.

When quark masses are taken into account, the fractions $x_a, x_b$ of hadron momenta carry by the incoming partons in hadron $A$ and hadron $B$ respectively become

$$\begin{aligned} x_a &= \sqrt{\tau}\left(\beta_i\cosh y_+ + \sinh y_+\right) \\ x_b &= \sqrt{\tau}\left(\beta_i\cosh y_+ - \sinh y_+\right). \end{aligned}$$

The integration range in equation (12) is determined by the cut $|\, y_1 \,| \leq 2.5, |\, y_2 \,| \leq 2.5$, which we impose on the rapidities of $W^-, W^+$, and by the conditions $\tau \leq x_a, x_b \leq 1$. The lower bound on $x_a, x_b$ is put in by hand as explained before. They have noticeable effects only on top quark luminosity.

We shall present our numerical results only for the diagonal cases, i.e., for $\lambda_1' = \lambda_1, \lambda_2' = \lambda_2$. These differential cross sections can be extracted, in principle, from measurements of the polar angle correlations of, say, the two charge leptons in the final state [19].

In our numerical work, we take $m_w = 81.8GeV, \sin^2\theta_w = 0.226, m_u = 3MeV, M_d = 5MeV, m_s = 150MeV, M_c = 1.5GeV, M_b = 5.5GeV, M_t = 70GeV, M_H = 400GeV$ and $\Gamma_H = 30GeV$. The quark mixing matrix is chosen to be orthegonal with $U_{ud} = 0.9750, U_{us} = 0.222, U_{ub} = 0.009539, U_{cd} = -0.221, U_{cs} = 0.9644, U_{cb} = 0.1455, U_{td} = 0.02311, U_{ts} = -0.144, U_{tb} = 0.9893$.

In Fig. 1, we compare the differential cross sections $[d\sigma/dMd\cos\theta^*]_{\lambda_1\lambda_2}$ at fixed $\theta^*$ for the cases when we choose fermion masses as above and when we set all fermion masses to zero at SSC energy and with the Standard Model couplings. The differences depend on $\theta^*$ and can be more than an order of magnitude in some cases. There is, however, very little difference in the (12) polarization channel[3] which dominates the polarization summed differential cross section at large $M$. In the (33) and (11)+(22) channels, we see the typical resonance effect due to the neutral scalars at 400 GeV. In general, with non-zero fermion masses, the differential cross sections tend to be more isotropic in $\theta^*$ distributions for all polarization channels.

Fig. 2 shows the effects of the form factors $f_8^Z, f_9^Z$ and $f_{10}^Z$ at Tevatron energy. We take each form factor to be 1.0 in turn. $f_8$ and $f_9$ affect (33) channel, $f_{10}$ affects (12) channel and $f_8$ also affects (11)+(22) channel. The differences from the Standard Model values are small except for the case $f_9 = 1.0$. Contributions from $f_i^S$ and $f_i^P$ are also small if they are of order one.

At SSC energy and with the rapidity cut we made, the angular dependence of the differential cross sections from $5^o$ to $90^o$ are rather flat except a dip at $90^o$ in the (11)+(22) channel. In Fig. 3 and Fig. 4 we show the effects of the nine form facts $f_8^Z, f_9^Z, f_{10}^Z, f_i^P, f_i^S$ when they take the value 1.0 in turn at SSC energy. The contributions from $f_8^Z$ and $f_9^Z$ to (33) channel and $f_{10}^Z$ to (12) channel are quite large and still increasing at 1 TeV. The contributions from $f_i^S$ and $f_i^P$ are similar in magnitude and in their dependence on $\theta^*$ and $M$. However, the resonance peak for $f_i^S$ due to Higgs coincides with that for the Standard Model. The resonance

[3] We do not distinguish which $W$ has polarization $\lambda = 1$. Similar comment holds for other polarization channels.

peak of $f_i^P$, in general, will be different from that of the Standard Model. In these graphs, we took all scalar masses to be 400 GeV.

In summary, we have given the formulas for the $W$ pair polarization density matrix taking the fermion masses into account. We also include in our formulas the effects of the most general neutral scalar and vector boson couplings to $W^+W^-$. We investigate, in particular, the numerical importance of the terms whose contributions to the differential cross sections are proportional to the incoming fermion masses. These terms have always been ignored in the previous discussions. This is certainly justifiable for electron machines and, to a certain extent, also true for the Tevatron unless luminosity is increased so that a study of different polarization channels of $W^+W^-$ is possible. However, the fact that they can be neglected also means that we would not be able to put any constraints on such terms through experiments done at these machines.

At SSC energy away from resonances, a value of order 0.1 for $f_8, f_9, f_{10}$ and $f_2^S, f_2^P$ will produce more than an order of magnitude deviations from the standard model predictions. For $f_{1,3}^S$ and $f_{1,3}^P$, they have to be of order 1 to produce the same effect. We have not taken into consideration the QCD background [14,20,27]. Moreover, we have tried to get some ideas of numerical importance of various form factors in a model independent way. For specific models, there may exist relations among the various form factors making the effects smaller in some cases.

We have studied only the diagonal elements of the polarization density matrix which may be extracted from the polar angle correlations of the final state fermions. Off-diagonal elements will contribute to the azimuthal angle correlations. A more detailed account of the present work and further studies will be presented elsewhere.

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## Figure Captions

Fig. 1: Differential cross sections for $W$ pair production at SSC energy with fermion masses set to zero (dotted) and with realistic fermion masses (solid). Polarization channels and $\theta^*$ are as indicated.

Fig. 2: Differential cross sections for $W$ pair production at Tevatron energy with Standard Model couplings (solid) and with $f_8^Z$ (dot-dashed (33) and dotted (11) + (22)), $f_9^Z$ (dotted, (33)) and $f_{10}^Z$ (dotted, (12)) set to 1.

Fig. 3: Differential cross sections for $W$ pair production SSC energy with Standard Model couplings (solid) and with $f_8^Z, f_9^Z$ (dashed); $f_1^S, f_2^S$ (dot-dashed); $f_1^P, f_2^P$ (dotted) set to 1. $f_9, f_2$ are plotted on the right hand graph.

Fig. 4: Same as Fig. 3, but for $f_{10}^Z, f_8^Z$ (dashed); $f_3^S, f_1^S$ (dot-dashed), $f_3^P, f_1^P$ (dotted) set to 1. $f_{10}, f_3$ are for (12) channel.

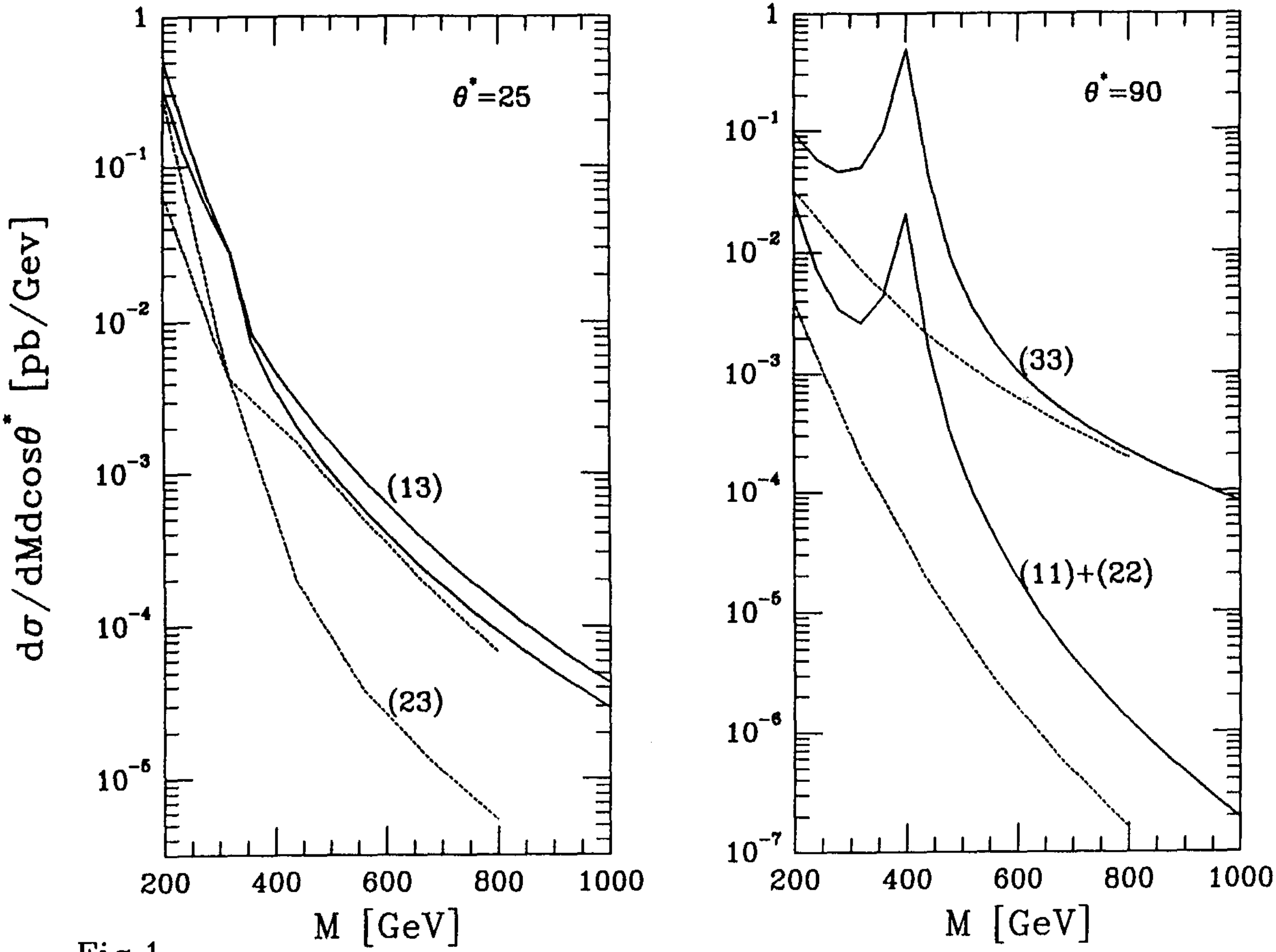
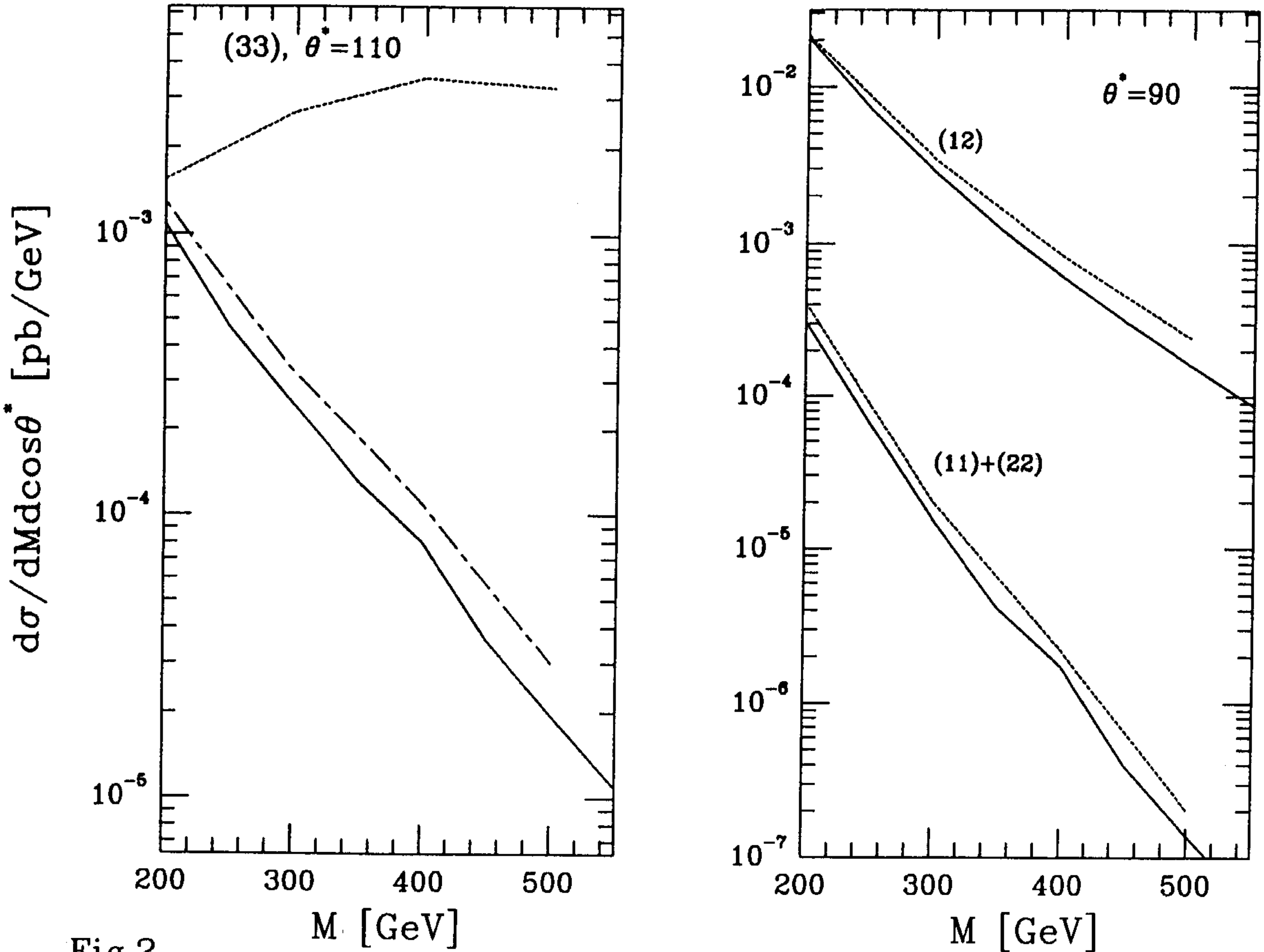
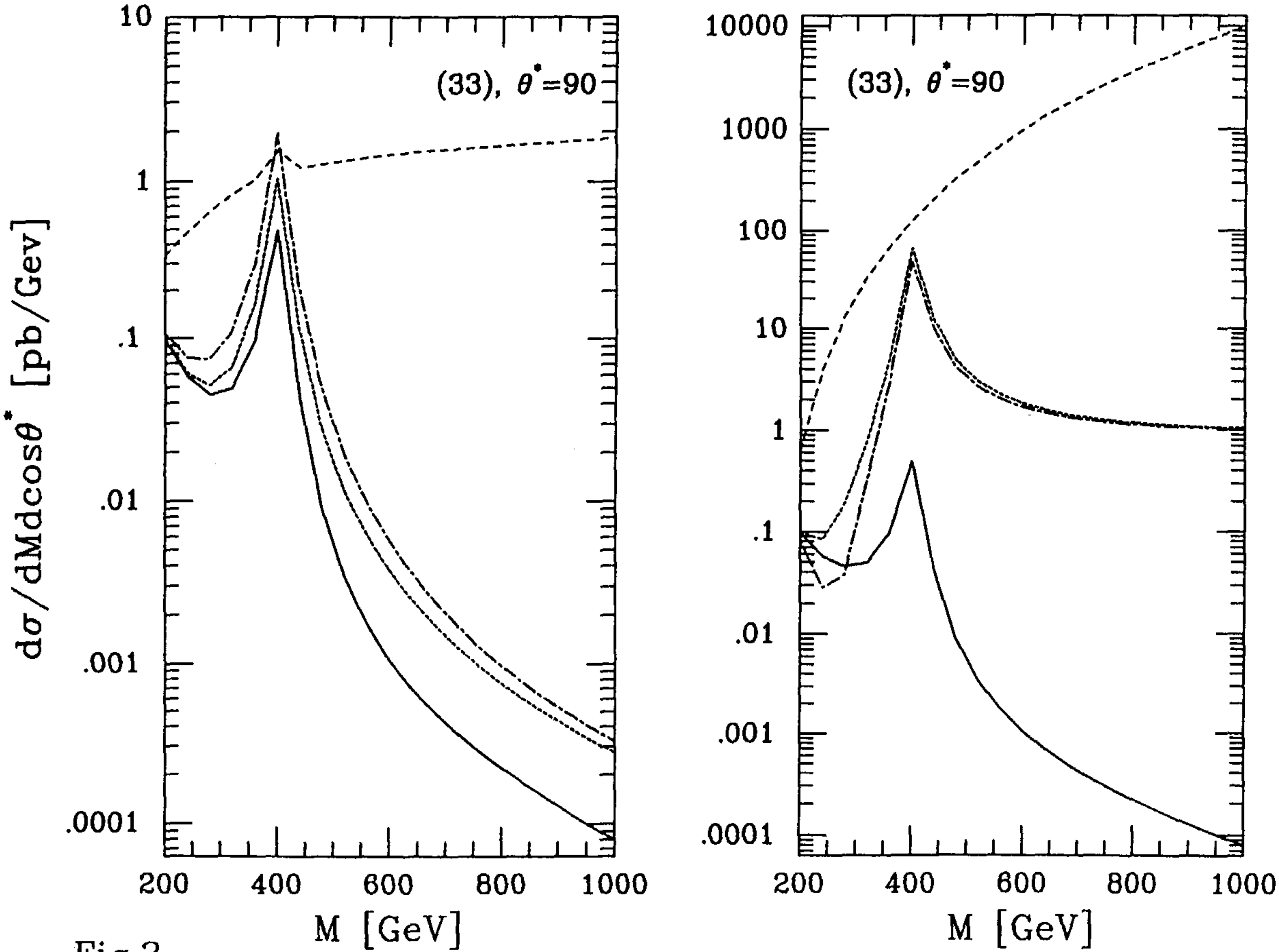

dσ/dMdcosθ* [pb/Gev]
1
10^-1
10^-2
10^-3
10^-4
10^-5
θ*=25
(13)
(23)
200
400
600
800
1000
M [GeV]
1
10^-1
10^-2
10^-3
10^-4
10^-5
10^-6
10^-7
θ*=90
(33)
(11)+(22)
200
400
600
800
1000
M [GeV]

Fig.1



dσ/dMdcosθ* [pb/GeV]
(33), θ*=110
10⁻³
10⁻⁴
10⁻⁵
200
300
400
500
M [GeV]
θ*=90
(12)
(11)+(22)
10⁻²
10⁻³
10⁻⁴
10⁻⁵
10⁻⁶
10⁻⁷
200
300
400
500
M [GeV]


Fig.2



dσ/dMdcosθ* [pb/Gev]
10
1
.1
.01
.001
.0001
200
400
600
800
1000
M [GeV]
(33), θ*=90
10000
1000
100
10
1
.1
.01
.001
.0001
200
400
600
800
1000
M [GeV]
(33), θ*=90


Fig.3

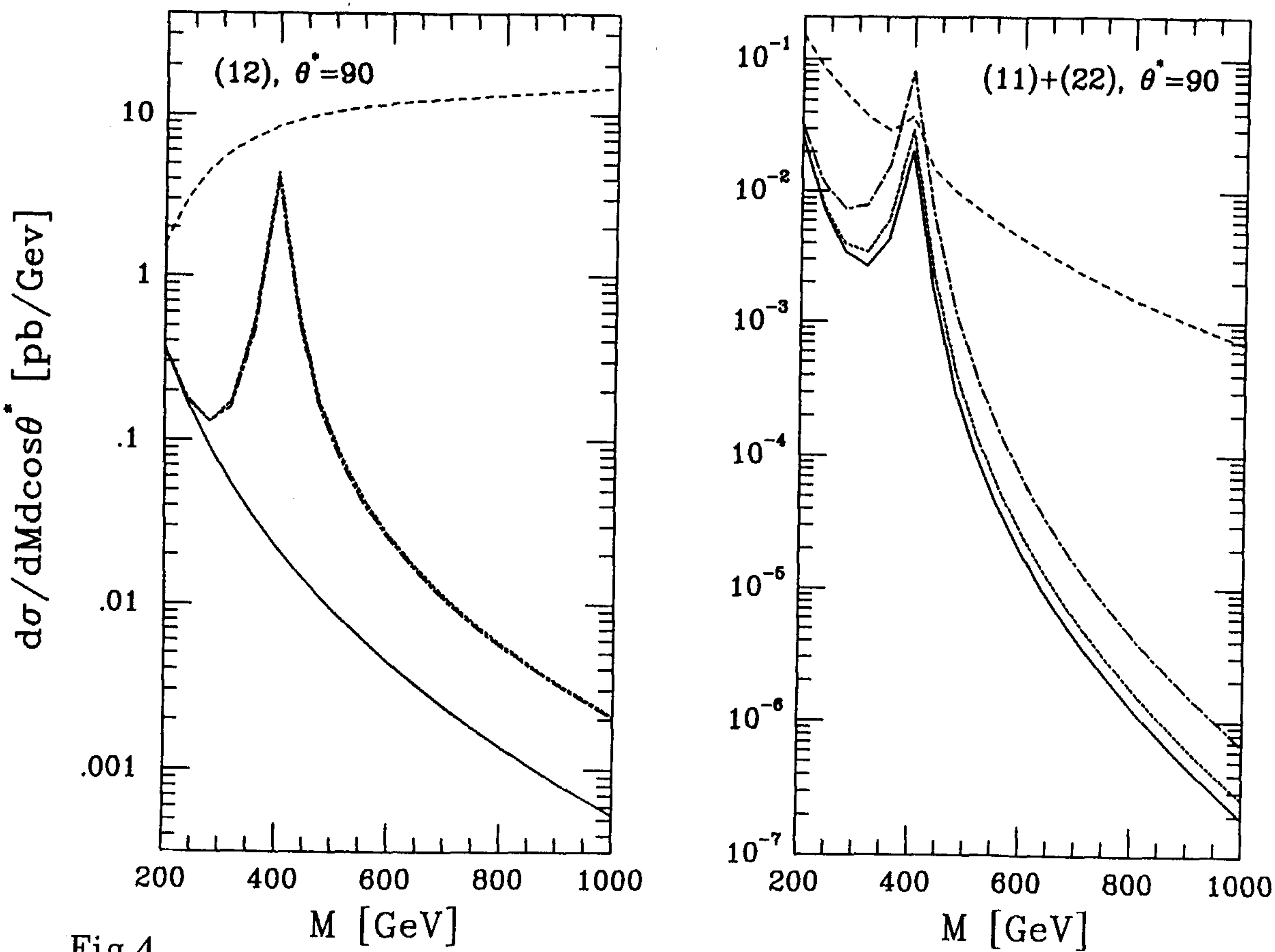

dσ/dMdcosθ* [pb/Gev]
(12), θ*=90
10
1
.1
.01
.001
200
400
600
800
1000
M [GeV]
(11)+(22), θ*=90
10^-1
10^-2
10^-3
10^-4
10^-5
10^-6
10^-7
M [GeV]

Fig.4